QCD-like Hidden Sector Models
without the Polonyi Problem

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ABSTRACT

QCD-like hidden sector models of supersymmetry breaking are considered which do not suffer from a cosmological problem due to the Polonyi field. Avoidance of a light gluino leads to introduction of quasi-symmetry – symmetry broken explicitly only through gravitational effects.

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1. INTRODUCTION

Nonabelian gauge theories are expected to provide a scale of supersymmetry (SUSY) breaking which eventually induces electroweak symmetry breaking at the weak scale. Although the simplest possibility is given by a pure Yang-Mills (YM) theory in the hidden sector (with singlets),[1] it suffers from a cosmological problem due to overproduction of light particles called the Polonyi field,[2] which is a gauge singlet[†] to gravitationally transmit the SUSY breaking in the hidden sector to the visible sector.

In this paper, we explore the next simplest candidate – vector-like nonabelian gauge theory – to break SUSY dynamically. Namely, matter fields in a real representation are introduced in the hidden sector. It turns out that this approach circumvents the Polonyi problem to give viable models for dynamical SUSY breaking.

2. QCD-LIKE INTERACTION IN THE HIDDEN SECTOR

Let us consider supersymmetric $SU(N_c)$ gauge theory with N_f pairs of hidden quark chiral superfields Q and $\overline{Q}$ in the fundamental representations N_c and $\overline{N}_c$, respectively. The hidden quarks are put to be massless.[‡] Otherwise, they would naturally have masses of order the gravitational scale $M \sim 10^{18}$GeV, which would make them decouple essentially to yield an effective theory similar to the pure YM case.

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† An elementary singlet is adequit, in particular, to produce sizable gaugino and Higgsino masses (see section 5).

‡ More presicely, the hidden quark mass term is excluded, which can be imposed by an axial symmetry given below.

We also introduce a singlet chiral superfield Z, whose lowest component is denoted by z. In contrast to the pure YM case, our QCD-like hidden sector contains a marginal interaction with the Polonyi superfield Z:

$$W_0 = \lambda_0 Z Q \overline{Q}, \tag{1}$$

which is crucial to avoid the Polonyi problem.[§] Here the self interactions of the field Z are excluded in accord with an axial $U(1)_A$ symmetry

$$Z \to e^{-2i\xi} Z, \quad Q \to e^{i\xi} Q, \quad \overline{Q} \to e^{i\xi}\overline{Q} \tag{2}$$

which removes the hidden quark mass term.

It is important that the global symmetry $U(1)_A$ has a hidden QCD anomaly, which allows a nonperturbative generation of an effective superpotential for Z violating $U(1)_A$. In fact, under the circumsatances that $\langle z \rangle \neq 0$, the hidden gauge and quark superfields may be integrated out to yield an effective superpotential[4] [¶]

$$W_1 = \lambda \Lambda^{3-n} Z^n; \quad n = \frac{N_f}{N_c}, \tag{3}$$

where Λ denotes the hidden QCD scale, which is supposed to satisfy $|\langle z \rangle| \ll \Lambda \ll M$. On the other hand, the Kähler potential then takes the following form:

$$K = ZZ^* - \frac{k}{2\Lambda^2}(ZZ^*)^2 + \cdots, \tag{4}$$

where k is real. Here the ellipsis denotes higher-order terms of ZZ^*, which are negligible in the subsequent analysis as far as $|\langle z \rangle| \ll \Lambda$. The dimensionless

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[§] Only M-suppressed interactions with Z are present in the pure YM case, which inevitablly causes the Polonyi problem.[3]

[¶] $N_f < 3N_c$ for asymptotic freedom. The form of this effective superpotential is understood by means of the non-anomolous $U(1)_R$ symmetry existing in the massless QCD with the superpotential (1).

constants λ and k are expected to be of order one. Note that the higher-dimensional terms in Eq. (4) are suppressed by Λ rather than M in contrast to the pure YM case.

We also introduce a constant term w to get

$$W = w + W_1 = w + \lambda \Lambda^{3-n} Z^n \tag{5}$$

as a total effective superpotential, where w and λ are chosen to be real. Though such a constant term may be induced by some dynamics outside the present hidden QCD, we simply regard it as an input parameter in this paper.

In the following discussion, we restrict ourselves to the case of $N_f = N_c$, in which the condition $0 \neq |\langle z \rangle| \ll \Lambda$ turns out to be consistent with the effective potential for z obtained in the next section.

3. Effective Potential for the Polonyi Field

The effective potential for the Polonyi field z in the presence of supergravity is given by

$$V = \exp\left(\frac{K}{M^2}\right) \left\{ \left(\frac{\partial^2 K}{\partial z \partial z^*}\right)^{-1} \left|\frac{\partial W}{\partial z} + \frac{\partial K}{\partial z}\frac{W}{M^2}\right|^2 - \frac{3|W|^2}{M^2} \right\}. \tag{6}$$

Requirement of vanishing cosmological constant on the vacuum with $|\langle z \rangle| \ll \Lambda \ll M$ implies that $|w|M^{-1} < \infty$ when $M \to \infty$. Then we obtain an approximate expression for the potential

$$\begin{aligned} V &\simeq \left(1 + \frac{2k}{\Lambda^2}zz^*\right) \left|\lambda\Lambda^2 + z^*\frac{w}{M^2}\right|^2 - \frac{3}{M^2}|w + \lambda\Lambda^2 z|^2 \\ &\simeq \lambda^2\Lambda^4 - \frac{3}{M^2}w^2 - \frac{4}{M^2}w\lambda\Lambda^2 x + 2k\lambda^2\Lambda^2(x^2 + y^2); \quad z = x + iy, \end{aligned} \tag{7}$$

where x and y are real and we suppose $k > 0$ from convexity of the effective

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potential in the limit $M \to \infty$. This results in the flat vacuum[*] with

$$w \simeq \frac{1}{\sqrt{3}} \lambda \Lambda^2 M, \quad \langle z \rangle \simeq \frac{\Lambda^2}{\sqrt{3}kM}, \tag{8}$$

which confirms the consistency $0 \neq |\langle z \rangle| \ll \Lambda$. This vacuum breaks SUSY with the breaking scale $M_S^2 \simeq \sqrt{3}|\langle W \rangle|M^{-1} \sim \Lambda^2$.

4. LOW-ENERGY CONTENTS IN THE HIDDEN SECTOR

The gravitino mass is given by

$$m_{3/2} \simeq \frac{|\langle W \rangle|}{M^2} \simeq \frac{\lambda \Lambda^2}{\sqrt{3}M}, \tag{9}$$

which characterizes the effective SUSY breaking scale in the visible sector and eventually the weak scale. Hence we set $\Lambda \sim 10^{10}\text{GeV}$ to get $m_{3/2} \sim 10^2\text{GeV}$.

The mass of the Polonyi field is of order Λ as seen in Eq. (7), which shows that the present model does not suffer from the Polonyi problem as desired.

The interaction (1) indicates that the hidden squarks acquire mass given by the F-component of Z, which is of order $M_S^2 \sim \Lambda^2$. Therefore the low-energy contents in the hidden sector are effectively described by non-supersymmetric QCD of hidden quarks (and 'gauginos' – see section 5) with mass $m = \lambda_0|\langle z \rangle| \sim 10^2\text{GeV}$. Thus there seems to exist hidden pions (and 'R-axion'[6]) with mass of order $\sqrt{m\Lambda} \sim 10^6\text{GeV}$, which may dominantly decay into gravitinos. This decay does not produce any cosmological problems if the reheating temperature after inflation is sufficiently low.[7]

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[*] Although it may not correspond to an absolute minimum of the potential, this vacuum is practically stable once a sufficiently large flat universe is formed.[5]

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5. Gaugino Mass and Quasi-symmetry

In order to give sizable masses to gauginos such as the gluino, we introduce higher-dimensional terms[8] of the form

$$\frac{1}{M} Z W_\alpha W^\alpha, \tag{10}$$

where W_α denote field-strength chiral superfields for gauge multiplets. The order of gaugino masses is then given by $M^{-1} M_S^2 \sim 10^2 \text{GeV}.^\star$

Though higher-dimensional terms are common in supergravity, the terms (10) pose a naturalness problem: the symmetry $\text{U}(1)_A$ defined by (2) removes the hidden quark mass term naturally, whereas it cannot help excluding (10) simultaneously.

Necessity of the terms (10) leads us to consider $\text{U}(1)_A$ as a quasi-symmetry – symmetry broken explicitly only through gravitational effects. The quasi-symmetry is supposed to be exact in the limit $M \to \infty$. Thus the presence of M-suppressed terms (10) is consistent with the notion of quasi-symmetry.† Moreover, it may be argued that a global symmetry is naturally a quasi-symmetry since its global charge 'dissipates' through black holes or wormholes[10] in quantum theory of gravitation.‡

6. Conclusion

We have considered QCD-like hidden sector models with the superpotential (1) on which the quasi-symmetry (2) is implimented. They turned out to suffer from

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$\star$ The Higgsino mass is also induced by a higher-dimensional term[9] of the form $M^{-1} Z^* H \overline{H}$ in the Kähler potential, where H and $\overline{H}$ denote doublet Higgs chiral superfields. Its order is given by $M^{-1} M_S^2$.

$\dagger$ The quasi-symmetry $\text{U}(1)_A$ allows nonrenormalizable interactions such as $M^{-1} Z^4$ in the superpotential, whose presence never affects our conclusion.

$\ddagger$ However, it is not clear to us whether breaking terms of the global symmetry are always suppressed by the gravitational scale M.

no cosmological problems and to give sizable gaugino masses. Thus we conclude that they constitute simple and viable models of SUSY breaking in the hidden sector.

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